Impact of spin-zero particle-photon interactions on light polarization in external magnetic fields

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Abstract

If the recent PVLAS results on polarization changes of a linearly polarized laser beam passing through a magnetic field are interpreted by an axion-like particle, it is almost certain that it is not a standard QCD axion. Considering this, we study the general effective interactions of photons with spin-zero particles without restricting the latter to be a pseudo-scalar or a scalar, i.e., a parity eigenstate. At the lowest order in effective field theory, there are two dimension-5 interactions, each of which has previously been treated separately for a pseudo-scalar or a scalar particle. By following the evolution in an external magnetic field of the system of spin-zero particles and photons, we compute the changes in light polarization and the transition probability for two experimental setups: one-way propagation and round-trip propagation. While the first may be relevant for astrophysical sources of spin-zero particles, the second applies to laboratory optical experiments like PVLAS. In the one-way propagation, interesting phenomena can occur for special configurations of polarization where, for instance, transition occurs but light polarization does not change. For the round-trip propagation, however, the standard results of polarization changes for a pseudoscalar or a scalar are only modified by a factor that depends on the relative strength of the two interactions.

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The neutral photon can interact with itself due to quantum effects. In QED this is summarized in the celebrated Euler-Heisenberg Lagrangian for weak external fields and low frequency photons [1]. Among physical effects of photon field self-interactions is the nonlinear optical phenomenon of a light beam in a magnetized vacuum [2, 3]. This latter effect has long been searched for in laboratory experiments without success [4] until recently by PLVAS, in which the rotation of the plane of polarization (dichroism) [5] and ellipticity (birefringence) [6] of a linearly polarized laser beam are observed after it traverses a transverse strong magnetic field. It is amusing that the observed rotation and ellipticity exceed the QED expectation by many orders of magnitude [2, 7, 8]. Without restricting to QED but allowing all possible photon field self-interactions, it has been shown that it is still not possible to accommodate the PVLAS results of rotation and ellipticity at the first non-trivial order in the low energy theory of photons [9]. This would mean that some ultra-light particles with a mass of order the laser frequency or lower and interacting with photon fields have to be invoked. The real and virtual production of these particles by laser in the external field induces the optical changes of the initial beam [10, 11, 12]. If the PVLAS results are confirmed, it would signal some new physics containing ultra-light particles.

The theoretically best motivated candidate for such a particle seems to be the axion, a pseudoscalar particle resulting from spontaneous breakdown of the Peccei-Quinn symmetry. In addition to offering a natural resolution to the strong CP problem, the axion also serves as an attractive candidate for dark matter. However, its mass and coupling to photons as determined by PVLAS are strongly excluded by helioscope experiments [13, 14] and astrophysical observations [15]. There are many discussions in the literature on this potential conflict, most appealing to hitherto unknown effects in the stellar environment [16]. An alternative to the axion or axion-like particles is the milli-charged particles with mass of order 0.1 eV and an electric charge of order $10^{-6}\ e$, interacting with photons as in QED [17]. These particles were shown some years ago to appear in para-photon models through kinetic mixing [18], and have recently been argued to appear naturally in string-based models [19]. They have also been utilized to circumvent the above mentioned conflict [20], and may have detectable effects in accelerator cavities [21] and reactor neutrino experiments [22]. Whether any of these suggestions is relevant will be examined in the new generation of experiments; in particular, axion-like particles can be best scrutinized in photon regeneration experiments [10, 4, 23, 24], one of which will start to operate soon by the ALPS group at DESY [25].

It is not the purpose of this work to add another proposal to relax the tension between the PVLAS and other experiments. Instead, our work is based on the observation [15] that, if the PVLAS results are interpreted in terms of an axion-like particle, it cannot be a QCD axion because the implied relation between mass and coupling is far away from the parameter region favored by PVLAS. In this circumstance there is no theoretical prejudice for a pseudoscalar or a scalar to explain the PVLAS results. In the effective theory at the low energy scale (~ 1 eV), the spin-zero particle probed by PVLAS could

well be one of no definite parity, i.e., a mixture of pseudoscalar and scalar. The portion of mixing depends on the details of some underlying fundamental theory that produces the ultra-light particle. Such particles could appear in extensions of the standard model as studied, for instance, in Ref.[26]. Then we should include in our effective theory all possible interactions of the particle with the photon field that have the same dimension. Note that there is no direct threat to experiments on parity and CP violation at much higher energy scales where our effective theory does not generally apply.

The effective theory at the lowest non-trivial order is defined by

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{1}{2}(\partial^\mu\varphi)^2 - \frac{1}{2}m^2\varphi^2 + \frac{1}{4}\lambda_+\varphi F^{\mu\nu}F_{\mu\nu} + \frac{1}{4}\lambda_-\varphi\tilde{F}^{\mu\nu}F_{\mu\nu}, \tag{1}$$

where φ is a spin-zero field, $F_{\mu\nu}$ the electromagnetic field tensor with dual $\tilde{F}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\alpha\beta}F_{\alpha\beta}$ ($\epsilon^{0123} = +1$), and $\lambda_\pm$ are two coupling constants. The separate cases with either $\lambda_+ = 0$ (φ being a pseudoscalar) or $\lambda_- = 0$ (a scalar) have been studied in the literature. The equations of motion (EoM's) are

$$\begin{aligned}(\partial^2 + m^2)\varphi &= \frac{1}{4}\lambda_+F^{\mu\nu}F_{\mu\nu} + \frac{1}{4}\lambda_-\tilde{F}^{\mu\nu}F_{\mu\nu}\\ \partial_\mu F^{\mu\nu} &= \partial_\mu\left(\lambda_+\varphi F^{\mu\nu} + \lambda_-\varphi\tilde{F}^{\mu\nu}\right)\end{aligned} \tag{2}$$

Since the background classical field will be much larger in strength than the quantum ones, the leading effects to quantum fields can be approximated by linearizing the EoM's with respect to them:

$$\begin{aligned}(\partial^2 + m^2)\varphi &= \frac{1}{2}\lambda_+f_{\mu\nu}F^{\mu\nu} + \frac{1}{2}\lambda_-f_{\mu\nu}\tilde{F}^{\mu\nu}\\ \partial_\mu f^{\mu\nu} &= \left(\lambda_+F^{\mu\nu} + \lambda_-\tilde{F}^{\mu\nu}\right)\partial_\mu\varphi\end{aligned} \tag{3}$$

where F (f) stands for the background (laser light) field tensor, and Maxwell equations have been used to simplify the second equation. Denoting $f_{\mu\nu} = \partial_\mu a_\nu - \partial_\nu a_\mu$ and working in the gauge $a^0 = \nabla\cdot\mathbf{a} = 0$, the EoM's in an external magnetic field $\mathbf{B}$ simplify to

$$\begin{aligned}(\partial_t^2 - \nabla^2 + m^2)\varphi &= \lambda_+(\nabla\times\mathbf{a})\cdot\mathbf{B} + \lambda_-(\partial_t\mathbf{a})\cdot\mathbf{B}\\ (\partial_t^2 - \nabla^2)\mathbf{a} &= \lambda_+(\nabla\varphi)\times\mathbf{B} - \lambda_-(\partial_t\varphi)\mathbf{B}\end{aligned} \tag{4}$$

We follow Ref.[12] to derive the evolution equation for a plane-wave light beam propagating in the z direction and perpendicularly to $\mathbf{B}$. We can choose $a^3 = 0$, and a^1, a^2, φ depend only on (t, z). Since for a temporally constant $\mathbf{B}$, the energies of the laser photon and the φ particle are equal, we can remove t-dependence by substitutions $\mathbf{a} \to \mathbf{a}e^{-i\omega t}$, $\varphi \to \varphi e^{-i\omega t}$, where ω is the angular frequency of the laser light. For propagation in the $+z$ direction, we can substitute on the left-hand side of equations, $(\omega^2 + \partial_z^2) = (\omega - i\partial_z)(\omega + i\partial_z) \to 2\omega(\omega + i\partial_z)$, and on the right, $\partial_z \to +i\omega$. The resultant

errors are higher order in small parameters not to be considered here. The EoM's in the matrix form are

$$\left(1 + i\omega^{-1}\partial_z + \Omega\right) \begin{pmatrix} a^1 \\ a^2 \\ \varphi \end{pmatrix} = 0, \quad \Omega = \begin{pmatrix} 0 & 0 & +i\delta_- \\ 0 & 0 & +i\delta_+ \\ -i\delta_- & -i\delta_+ & -\delta_0 \end{pmatrix}, \tag{5}$$

where

$$\delta_- = \frac{1}{2\omega}(\lambda_- B_1 - \lambda_+ B_2), \ \delta_+ = \frac{1}{2\omega}(\lambda_- B_2 + \lambda_+ B_1), \ \delta_0 = \frac{m^2}{2\omega^2} \tag{6}$$

Without loss of generality, we choose $\mathbf{B} = |\mathbf{B}|\hat{x}$, then $\delta_\pm = (2\omega)^{-1}\lambda_\pm|\mathbf{B}|$. In this coordinate system, $a^{1,2}$ are respectively the components parallel and perpendicular to $\mathbf{B}$, thus denoted below as $a_{\parallel,\perp}$.

To find the eigenmodes of propagation, we diagonalize the matrix Ω in steps, with the result

$$\begin{aligned} \Omega_{\rm diag} &= U^{-1}\Omega U = \delta_0 \ \mathrm{diag}\left(0, \epsilon^2, -(1+\epsilon^2)\right) + \delta_0 O(\epsilon^4), && (7) \\ U &= E R_a(\theta_\lambda) R_{a\varphi}(\epsilon) = E \begin{pmatrix} c_\lambda & s_\lambda c_\epsilon & -s_\lambda s_\epsilon \\ -s_\lambda & c_\lambda c_\epsilon & -c_\lambda s_\epsilon \\ 0 & s_\epsilon & c_\epsilon \end{pmatrix} && (8) \end{aligned}$$

where $E = \mathrm{diag}(i, i, 1)$ removes the $\pm i$ factors in Ω, $R_a(\theta_\lambda)$ rotates the two components $\delta_\pm$ into one,

$$\delta = \sqrt{\delta_+^2 + \delta_-^2} = \frac{|\mathbf{B}|}{2\omega}\sqrt{\lambda_+^2 + \lambda_-^2} \tag{9}$$

which in turn mixes with φ. This last mixing is then diagonalized by $R_{a\varphi}(\epsilon)$. The following short-cuts

$$s_\lambda = \sin\theta_\lambda = \frac{\lambda_-}{\sqrt{\lambda_+^2 + \lambda_-^2}}, \ c_\lambda = \cos\theta_\lambda = \frac{\lambda_+}{\sqrt{\lambda_+^2 + \lambda_-^2}} \tag{10}$$

$$s_\epsilon = \sin\epsilon, \ c_\epsilon = \cos\epsilon, \ \tan 2\epsilon = \frac{2\delta}{\delta_0} = \frac{2\omega|\mathbf{B}|}{m^2}\sqrt{\lambda_+^2 + \lambda_-^2} \tag{11}$$

have been used. For PVLAS one has $\omega = \frac{2\pi}{\lambda} \approx 1.17$ eV, $|\mathbf{B}| = 5$ Tesla ≈ 976.68 eV^2. PVLAS found that one would want $\lambda_\pm \sim (4 \times 10^5 \text{ GeV})^{-1}$, $m \sim 10^{-3}$ eV to explain the results. Thus the above mixing is very small, $\epsilon \sim 3 \times 10^{-6}$.

A state that enters the $\mathbf{B}$ field at z_0, $\Psi(z_0) = (a_\parallel(z_0), a_\perp(z_0), \varphi(z_0))^T$ will evolve into the state, $\Psi(z) = (a_\parallel(z), a_\perp(z), \varphi(z))^T$, after traversing a distance $z - z_0 = L \geq 0$ in the $+z$ direction:

$$\Psi(z_0 + L) = e^{i\omega L} V_+(L)\Psi(z_0), \ V_+(L) = U e^{i\omega L \Omega_{\rm diag}} U^{-1} \tag{12}$$

Expansion in ϵ yields

$$V_+(L) = V_0(L) + \epsilon V_1(L) + \epsilon^2 V_2(L) + O(\epsilon^3) \tag{13}$$

where

$$\begin{aligned}
V_0(L) &= \mathrm{diag}(1,1,e^{-i\zeta}),\ \zeta = \delta_0\omega L & (14)\\
V_1(L) &= (1-e^{-i\zeta})\begin{pmatrix} & & is_\lambda \\ & & ic_\lambda \\ -is_\lambda & -ic_\lambda & 0\end{pmatrix} & (15)\\
V_2(L) &= (-1+i\zeta+e^{-i\zeta})\begin{pmatrix} s_\lambda^2 & c_\lambda s_\lambda & \\ c_\lambda s_\lambda & c_\lambda^2 & \\ & & 0\end{pmatrix}\\
&+ (1-e^{-i\zeta}(1+i\zeta))\mathrm{diag}(0,0,1) & (16)
\end{aligned}$$

For $\lambda_+ = 0$, we have $c_\lambda = 0,\ s_\lambda = \mathrm{sign}(\lambda_-)$ and $s_\lambda\epsilon = \omega|\mathbf{B}|\lambda_- m^{-2}$. This is the case elaborated upon in Ref.[12]. The difference to that reference is that the $\pm i$ factors multiplying $s_\lambda,\ c_\lambda$ in $V_1(z)$ (which amount to a re-phasing of the photon state) were dropped there and that their ζ seems to be $-\zeta$ for propagation in the $+z$ direction.

For propagation in the $-z$ direction, the linearized EoM's can be obtained from eqn. (5) by $\partial_z \to -\partial_z$ and $\lambda_+ \to -\lambda_+$. The latter replacement amounts to $c_\lambda \to -c_\lambda$ in the matrices $U,\ V$ for our choice $\mathbf{B} = |\mathbf{B}|\hat{x}$. The former has no effect as long as the phase retardation is measured by the distance $L \geq 0$. Thus, a state, $\Psi(z_0)$, entering the $\mathbf{B}$ field at z_0 will evolve into the following one, after traversing a distance $z_0 - z = L$ in the $-z$ direction:

$$\Psi(z_0 - L) = e^{i\omega L}V_-(L)\Psi(z_0),\ V_-(L) = V_+(L)|_{c_\lambda\to -c_\lambda} \tag{17}$$

Now we consider the optical effects of the evolution equation. Suppose a beam of light propagates in the $+z$ direction and enters a transverse $\mathbf{B} = \mathbf{B}\hat{x}$ field. If the initial beam is linearly polarized at an angle θ measured counter-clockwise in the (xy) plane with respect to the field, it will evolve into the following state after traversing a distance L in the field:

$$\begin{pmatrix} a_\| \\ a_\perp \\ \varphi \end{pmatrix}(L) = e^{i\omega L}V_+(L)\begin{pmatrix} c_\theta \\ s_\theta \\ 0\end{pmatrix} \equiv e^{i\omega L}\begin{pmatrix} \eta\cos(\theta+\Delta\theta)e^{i\phi_\|} \\ \eta\sin(\theta+\Delta\theta)e^{i\phi_\perp} \\ \rho e^{i\sigma}\end{pmatrix} \tag{18}$$

where, to $O(\epsilon^2)$,

$$\begin{aligned}
\eta\cos(\theta+\Delta\theta) &= \cos\theta - \epsilon^2(1-\cos\zeta)\sin(\theta_\lambda+\theta)\sin\theta_\lambda\\
\eta\sin(\theta+\Delta\theta) &= \sin\theta - \epsilon^2(1-\cos\zeta)\sin(\theta_\lambda+\theta)\cos\theta_\lambda\\
\phi_\| &= \epsilon^2\sin(\theta_\lambda+\theta)\frac{\sin\theta_\lambda}{\cos\theta}(\zeta-\sin\zeta)\\
\phi_\perp &= \epsilon^2\sin(\theta_\lambda+\theta)\frac{\cos\theta_\lambda}{\sin\theta}(\zeta-\sin\zeta)\\
\rho e^{i\sigma} &= 2\epsilon\sin(\theta_\lambda+\theta)\sin\frac{\zeta}{2}\exp\left(-\frac{i}{2}\zeta\right)
\end{aligned} \tag{19}$$

Conservation of probability requires $\eta^2 + \rho^2 = 1$ as can be checked explicitly.

The polarization of the initial beam has been rotated at $O(\epsilon^2)$ by,

$$\Delta\theta = -\epsilon^2 \sin^2 \frac{\zeta}{2} \sin 2(\theta_\lambda + \theta) \tag{20}$$

The relative phase shift of the parallel and perpendicular components of the light beam results in the ellipticity measured by $\tan\chi$ where $\chi \in [-\frac{\pi}{4}, \frac{\pi}{4}]$ is determined by $\sin 2\chi = \sin 2(\theta + \Delta\theta) \sin(\phi_\perp - \phi_\parallel)$. Thus, at $O(\epsilon^2)$, we have

$$\tan\chi \approx \chi \approx \frac{1}{2}\epsilon^2(\zeta - \sin\zeta) \sin 2(\theta_\lambda + \theta) \tag{21}$$

Although the relative phase shift has to be treated carefully at the angles with $\sin 2\theta = 0$, the above ellipticity is always well defined. The production probability of φ is

$$P[(\gamma \text{ at } \theta) \to \varphi] = |\rho|^2 = 4\epsilon^2 \sin^2(\theta_\lambda + \theta) \sin^2 \frac{\zeta}{2} \tag{22}$$

The familiar results for a pure pseudoscalar or scalar are recovered by putting $c_\lambda = 0,\ |s_\lambda| = 1$ or $s_\lambda = 0,\ |c_\lambda| = 1$ respectively.

In a photon regeneration or shinning-light-through-walls experiment like ALPS, φ particles would be produced by laser in the production zone, then penetrate a wall that blocks the laser, and enter into the detection zone. This beam of φ particles will evolve into the state

$$e^{i\omega L}\left(i\epsilon s_\lambda(1 - e^{-i\zeta}), i\epsilon c_\lambda(1 - e^{-i\zeta}), e^{-i\zeta} + \epsilon^2[1 - e^{-i\zeta}(1 + i\zeta)]\right)^T \tag{23}$$

The probability to produce a photon is thus

$$P[\varphi \to \text{photon}] = 4\epsilon^2 \sin^2 \frac{\zeta}{2} \tag{24}$$

The produced photon is linearly polarized at the angle θ to the magnetic field, determined by $\tan\theta \tan\theta_\lambda = 1$. It would thus be possible to extract the parity property of φ if the photon's polarization could be measured.

A few comments concerning the above results are in order. When $\sin(\theta_\lambda + \theta) = 0$, i.e., $\tan\theta = -\lambda_-/\lambda_+$, there are no optical effects and a photon with this particular polarization relative to $\mathbf{B}$ cannot be converted into a φ, although an existing φ can still be converted into a photon with an orthogonal polarization. When $\cos(\theta_\lambda + \theta) = 0$, i.e., $\tan\theta = \lambda_+/\lambda_-$, there are no optical effects due to equal attenuation and phase retardation of the two orthogonal polarizations, but the photon to φ transition probability reaches its maximum which is equal to that of the inverse transition. Finally, for $\zeta = \frac{m^2 L}{2\omega} \ll 1$, the two optical quantities induced by a spin-zero particle are simply related by

$$\frac{\Delta\theta}{\chi} = -\frac{3}{\zeta} = -\frac{6\omega}{m^2 L} \tag{25}$$

While the above results may be relevant for astrophysical sources of φ particles, they are insufficient for laboratory optical experiments where the laser light is reflected back

and forth in a magnetic field for a larger gain of the signal. Due to parity violating interactions, the changes of polarization in opposite directions are not equal, and can cancel partially for a round trip. The detailed information on the evolution of a state $\Psi(0)$ can be obtained as follows for a round trip in the $\mathbf{B}$ field. Suppose it propagates first for a distance L in the $+z$ direction and perpendicularly to $\mathbf{B}$, gets reflected by a high-reflectivity mirror, then propagates in the $-z$ direction for the same distance and is reflected back to its starting point. The state becomes

$$\Psi_{\rm round} = (Re^{i\omega L}V_-)(Re^{i\omega L}V_+)\Psi(0) \tag{26}$$

where $R = \mathrm{diag}(1, 1, 0)$ expresses the fact that the produced φ particles are not reflected but penetrate the mirror to disappear. The evolution is thus not coherent at the mirrors. Irrelevant global phases at reflection have been neglected in the above.

In all optical experiments the light is reflected many times. Since the signal is extremely small for one passage in the magnetic field, the final state can be well approximated by that after N times of round trips:

$$\Psi_{N,\rm round} = e^{i2\omega NL}(RV_-RV_+)^N\Psi(0) \tag{27}$$

where N is roughly half the total number of passage of light in the field. Using

$$(RV_-RV_+)^N = \begin{pmatrix} C^N & 0 & C^{N-1}X \\ 0 & S^N & S^{N-1}Y \\ 0 & 0 & 0 \end{pmatrix} \tag{28}$$

where, to $O(\epsilon^2)$,

$$\begin{aligned} C &= 1 + 2\epsilon^2 s_\lambda^2(-1 + i\zeta + e^{-i\zeta}) \\ &\approx [1 - 4\epsilon^2 s_\lambda^2 \sin^2\tfrac{\zeta}{2}]e^{i2s_\lambda^2\epsilon^2(\zeta - \sin\zeta)} \\ S &= 1 + 2\epsilon^2 c_\lambda^2(-1 + i\zeta + e^{-i\zeta}) \\ &\approx [1 - 4\epsilon^2 c_\lambda^2 \sin^2\tfrac{\zeta}{2}]e^{i2c_\lambda^2\epsilon^2(\zeta - \sin\zeta)} \\ X &= i\epsilon s_\lambda(1 - e^{-i\zeta}) \\ Y &= i\epsilon c_\lambda(1 - e^{-i\zeta}) \end{aligned} \tag{29}$$

an initial laser beam of $\Psi(0) = (c_\theta, s_\theta, 0)^T$ will evolve into the state

$$\Psi_{N,\rm round} = e^{i2\omega NL}\begin{pmatrix} \left[1 - 4N\epsilon^2 s_\lambda^2 \sin^2\frac{\zeta}{2}\right] e^{i2N\epsilon^2 s_\lambda^2(\zeta - \sin\zeta)} \\ \left[1 - 4N\epsilon^2 c_\lambda^2 \sin^2\frac{\zeta}{2}\right] e^{i2N\epsilon^2 c_\lambda^2(\zeta - \sin\zeta)} \end{pmatrix} \tag{30}$$

for $N\epsilon^2 \ll 1$. The induced rotation and ellipticity are

$$\Delta\theta = -2N\epsilon^2 \sin^2\frac{\zeta}{2}\sin 2\theta\cos 2\theta_\lambda \tag{31}$$

$$\chi = N\epsilon^2(\zeta - \sin\zeta)\sin 2\theta\cos 2\theta_\lambda \tag{32}$$

which are modified by a factor $\cos 2\theta_\lambda$ from the standard results for a particle of definite parity. This simple modification means that the potential tension between the PVLAS

result on rotation (pointing to a pseudoscalar) [5] and its preliminary result on ellipticity (favoring a scalar) [6] is not relaxed in the parity non-conserving case. When the two interactions are of the same strength, i.e., $\cos 2\theta_\lambda = 0$, no net optical effects remain as intuitively expected, although the photon to φ transition generally occurs. It is also clear that a reflection symmetric experiment like PVLAS cannot tell a pseudoscalar or scalar particle from one of no definite parity.

If our aim is to work out the rotation and ellipticity at the first non-trivial order in ϵ, there is a simpler way to proceed. The amount collected on the return trip can be obtained from that on the forward trip by $c_\lambda \to -c_\lambda$, which is equivalent to $\theta_\lambda \to -\theta_\lambda$ in eqns. (20,21). Adding the amount for a round trip yields the result shown in eqn (32) for $N = 1$.

To summarize, we have studied the general electromagnetic couplings of a neutral, spin-zero particle in an effective field theory at the energy scale of order eV. Our consideration was based on the observation that, if the PVLAS results are explained in terms of an axion-like particle, the latter cannot be a QCD-like axion. In this circumstance, there is no theoretically strong preference for a pseudoscalar or scalar particle; instead, it could be a particle of no definite parity. We have considered the effects of such a particle on the evolution of a light beam propagating in a transverse magnetic field, and calculated the changes in the light polarization for two experimental set-ups. For a parity asymmetric set-up, e.g., with light propagating in one direction, interesting phenomena can occur for special polarization configurations where neither rotation nor ellipticity is induced although particle transitions still take place. Those configurations are determined by the relative strength of the interactions. In a parity symmetric set-up like PVLAS, however, optical changes are simply modified by a common factor, compared to the pure-parity case. The factor is again fixed by the relative strength of the interactions. Thus, if the results in such an experiment can be explained in terms of a particle with definite parity, they can be equally well explained in terms of a particle of no definite parity by adjusting the factor. Nevertheless, the parity property could be studied in a photon regeneration experiment if the polarization of regenerated photons could be measured.

In optical experiments like BFRT and PVLAS, a very small amount of gas is usually introduced into the cavity for calibration purposes. This amounts to introducing new terms in our Ω shown in eqn.(5) to replace the diagonal zeros (due to Cotton-Mouton effect) and off-diagonal ones (for Faraday effect in the presence of a magnetic field component along the propagation direction). This more general mixing case is also relevant for phenomena in the stellar environment. The interplay with the simple mixing scheme discussed in this work, including possible modifications in pressure, magnetic field and path length dependencies, deserves further study to which we hope to come back soon.

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